Nonthermal CP violation in soft leptogenesis

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Soft leptogenesis is a mechanism which generates the matter-antimatter asymmetry of the Universe via the out-of-equilibrium decays of heavy sneutrinos in which soft supersymmetry breaking terms play two important roles: they provide the required CP violation and give rise to the mass splitting between otherwise degenerate sneutrino mass eigenstates within a single generation. This mechanism is interesting because it can be successful at the lower temperature regime $T \lesssim 10^9$ GeV in which the conflict with the overproduction of gravitinos can possibly be avoided. In earlier works the leading CP violation is found to be nonzero only if finite temperature effects are included. By considering generic soft trilinear couplings, we find two interesting consequences: (1) the leading CP violation can be nonzero even at zero temperature realizing nonthermal CP violation, and (2) the CP violation is sufficient even far away from the resonant regime allowing soft supersymmetry breaking parameters to assume natural values at around the TeV scale. We discuss phenomenological constraints on such scenarios and conclude that the contributions to charged lepton flavor violating processes are close to the sensitivities of present and future experiments.

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I. INTRODUCTION

Leptogenesis [1] is an attractive mechanism for generating the observed matter-antimatter asymmetry of the Universe wherein one first creates an asymmetry in the lepton sector which, in turn, induces an asymmetry in the baryon sector via anomalous $B+L$ violating interactions. In standard type-I seesaw supersymmetric leptogenesis [2–5] involving the out-of-equilibrium decays of heavy neutrinos and sneutrinos, the CP violation required to generate the lepton number asymmetry comes from the neutrino Yukawa couplings. This scenario, with hierarchical right-handed neutrinos (RHNs), faces a conflict as successful leptogenesis requires the mass of the lightest RHN to be at least 10^9 GeV [6] while the simplest resolution of the gravitino problem [7, 8] requires the reheating temperature after inflation to be less than 10^{6-9} GeV depending on the gravitino mass [9].[1]

One may avoid this conflict by incorporating new elements in leptogenesis. In models of soft leptogenesis [13, 14] (for a recent review, see Ref. [15]) CP violation comes from soft supersymmetry (SUSY) breaking terms (here onwards we will simply refer to them as soft terms) with soft parameters assumed to be at the $m_{\rm SUSY} \sim$ TeV scale; i.e., we still hope SUSY is responsible for stabilizing the hierarchy between the weak and grand unification scales. One interesting feature is that soft leptogenesis can proceed even with one generation of the RHN chiral superfield.[2] Essentially, the heavy sneutrino $\widetilde{N}$ and antisneutrino $\widetilde{N}^*$ from the same chiral supermultiplet will mix due to the presence of the soft terms. The decays of the mixed mass eigenstates violate both CP and lepton number and generate a matter-antimatter asymmetry. Although the CP violation is suppressed by powers of $m_{\rm SUSY}/M \ll 1$ with M the scale of the lightest RHN, the mass splitting between these otherwise degenerate sneutrino mass eigenstates is also proportional to $m_{\rm SUSY}/M$. Crucially, this small splitting also results in enhancement of the CP violation from mixing. Because of the suppression factor $m_{\rm SUSY}/M$ in the CP violation, one cannot have very large M. Estimating the leading CP parameter as $\epsilon \sim m_{\rm SUSY}/M$ and that successful leptogenesis generically requires $\epsilon \gtrsim 10^{-6}$, we obtain $M \lesssim 10^9$ GeV assuming $m_{\rm SUSY}$ at the TeV scale. Hence soft leptogenesis occurs in the regime where the conflict with the bound on the reheating temperature from gravitino overproduction can be mitigated or even avoided.

In the original proposals of Refs. [13, 14], the authors showed that in the scenario of $\widetilde{N} - \widetilde{N}^*$ mixing, the leading CP violation in decays to fermions and scalars have opposite signs and cancel each other at the order $\mathcal{O}\left(m_{\rm SUSY}/M\right)$ at zero temperature $T=0$. They further showed that once finite temperature effects are taken into account, this cancellation is partially lifted, i.e. one obtains an asymmetry proportional to a factor $[c_F(T) - c_B(T)]$, where $c_{F,B}(T)$ are phase space and statistical factors associated with fermion and boson final states, and where the contributions do not completely cancel each other at finite temperature. Working under the assumption of proportionality of soft trilinear couplings $A_\alpha = AY_\alpha$ where the Y_α's are the neutrino Yukawa couplings and α the lepton flavor index, they showed that the resulting CP violation is of the order of $\mathcal{O}\left(m_{\rm SUSY}/M\right)$ at the *resonance* which, however, requires an unconventionally small soft bilinear coupling $B \ll m_{\rm SUSY}$. Away from the resonance, the CP violation is of $\mathcal{O}\left(Y_\alpha^2\right)$ and, hence, too suppressed for successful leptogenesis. On the other hand, assuming generic A couplings, Ref. [16] showed that successful leptogenesis can be obtained with $B \sim m_{\rm SUSY}$ away from the resonant regime.

Later in Ref. [17] it was argued that direct CP violation, i.e., from vertex corrections, due to gaugino exchange in the loop, survives at the order $\mathcal{O}\left(m_{\rm SUSY}^2/M^2\right)$ at $T=0$. Since the neutrino Yukawa coupling is replaced by the gauge coupling in the CP violation parameter, a large CP violation can be obtained for M at the TeV scale. Further study in Ref. [18], however, showed that in fact in this scenario, the cancellation still holds up to $\mathcal{O}\left(m_{\rm SUSY}^2/M^2\right)$ at $T=0$, and it was concluded that finite temperature effects are necessary to prevent the cancellation. The cancellation is consistent with the result obtained in Ref. [19] which states that to have a nonvanishing total CP violation there should be lepton number violation to the right of the "cut" in the loop diagram, and this requirement is not fulfilled in these cases. More recently, in Ref. [20] it was shown that if finite temperature effects are taken into account *consistently*, the cancellation of direct CP violation from the gaugino contribution still holds even at $T \neq 0$.

In fact, in soft leptogenesis at finite temperature, the partial cancellation in the resulting lepton and slepton number density asymmetries sourced by CP violation from mixing and the complete cancellation in the case of the gaugino vertex correction [20] only hold under the assumption of equilibration between the chemical potentials of leptons and sleptons (superequilibration) which is valid below $T \lesssim 10^8$ GeV for $m_{\rm SUSY} \sim$ TeV [5]. As shown in Ref. [21], in the nonsuperequilibration regime, the partial cancellation between lepton and slepton number density asymmetries in the mixing scenario is avoided, resulting in an enhanced efficiency for soft leptogenesis. However, for reasons given later, we shall below consider mixing and vertex scenarios in the superequilibration regime (and also find a case where the lepton and slepton number density asymmetries do not partially cancel each other). On the other hand, considering $M \gtrsim 10^8$ GeV and $m_{\rm SUSY} \sim 1$ TeV, the CP violating parameter from the gaugino contribution in the

[1] See Refs. [10–12] for another resolution of the gravitino problem due to delayed thermalization of the Universe after inflation.

[2] In a realistic model, we need at least two RHNs to accommodate neutrino oscillations. Assuming RHNs to be hierarchical, soft leptogenesis only depends on the parameters related to the lightest RHN and decouples from the parameters related to heavier RHNs.

nonsuperequilibration regime is $\epsilon \sim 10^{-1} m_{\rm SUSY}^2/M^2 \lesssim 10^{-11}$ and, hence, is too small for successful leptogenesis. Therefore processes involving gauginos will not be considered further in this work.

In this article, we revisit soft leptogenesis by relaxing the assumption of the proportionality of the A couplings. In Sec. II, we review the Lagrangian for soft leptogenesis with generic A_α terms and spell out the constraints from out-of-equilibrium decays of heavy sneutrinos and the cosmological bound on the sum of neutrino masses. In Sec. III, we obtain the CP violating parameter for both the self-energy corrections (mixing) and vertex corrections. We show that generic A_α couplings give rise to two interesting consequences: (1) the leading CP violation can be nonzero even when thermal corrections are neglected implying a *nonthermal* CP violation, and (2) the mixing CP violation away from the resonance is of the order of $\mathcal{O}(Y_\alpha)$ and, hence, can be large enough for leptogenesis. Because of the small mass splitting, the mixing CP violation always dominates over the vertex CP violation even far away from the resonant regime. In Ref. [22], it was shown that with $A_\alpha = AY_\alpha$, soft leptogenesis gives negligible contributions to the electric dipole moment of charged leptons and charged lepton flavor violating processes. In Sec. IV, we repeat the exercise and show that with generic A_α couplings, the contributions to charged lepton flavor violating processes are close to the sensitivities of present and future experiments. Finally, in Sec. V, we conclude. This article is completed with two appendixes. In Appendix A, we discuss the inclusion of thermal effects under the assumption of decaying heavy sneutrinos at rest. In Appendix B, we review the two specific scenarios of A_α discussed in Ref. [16] and discuss an interesting point missed by Ref. [16] which actually allows for nonzero leading CP violation at zero temperature.

II. THE LAGRANGIAN

The superpotential for the type-I seesaw is given by

$$W_N = \frac{1}{2} M_i \hat{N}_i^c \hat{N}_i^c + Y_{i\alpha} \hat{N}_i^c \hat{\ell}_\alpha \hat{H}_u, \tag{1}$$

where $\hat{N}_i^c$, $\hat{\ell}_\alpha$ and $\hat{H}_u$ denote, respectively, the chiral superfields of the RHNs, the lepton doublet and the up-type Higgs doublet, and i and α are the RHN family and lepton flavor indices, respectively. The $SU(2)_L$ contraction between $\hat{\ell}_\alpha$ and $\hat{H}_u$ is left implicit. In the following, we will assume that the RHNs are hierarchical such that only the lightest RHN N_1 is relevant for soft leptogenesis. Henceforth, we will drop the family index of RHN, for example, $N \equiv N_1$ and $Y_\alpha \equiv Y_{1\alpha}$. The corresponding soft terms are

$$-\mathcal{L}_{\rm soft} = \widetilde{M}^2 \tilde{N}^* \tilde{N} + \left(\frac{1}{2} BM \widetilde{N}\widetilde{N} + A_\alpha \widetilde{N} \widetilde{\ell}_\alpha H_u + \text{H.c.} \right). \tag{2}$$

The mass and interaction terms involving the sneutrino $\widetilde{N}$ from W_N are given by

$$-\mathcal{L}_{\widetilde{N}} = |M|^2 \widetilde{N}^* \tilde{N} + \left(M^* Y_\alpha \widetilde{N}^* \widetilde{\ell}_\alpha H_u + Y_\alpha \overline{\widetilde{H}_u^c} P_L \ell_\alpha \widetilde{N} + \text{H.c.} \right), \tag{3}$$

where $P_{L,R} = \frac{1}{2}(1 \mp \gamma_5)$. Through field redefinitions, it can be shown that the three physical phases are

$$\Phi_\alpha = \arg \left(A_\alpha Y_\alpha^* B^* \right). \tag{4}$$

Without loss of generality, the phases can be assigned to A_α and all other parameters will be taken real and positive. We would like to stress that we do not assume the proportionality of A_α to the neutrino Yukawa couplings ($A_\alpha = AY_\alpha$) as has been done in Refs. [13, 14, 23] where there is only one physical phase $\Phi = \arg(AB^*)$. As we will show in Sec. III, by considering generic A_α couplings, the CP violation can be nonvanishing even at zero temperature.

Because of the bilinear B term, $\widetilde{N}$ and $\widetilde{N}^*$ mix to form mass eigenstates

$$\begin{aligned} \widetilde{N}_+ &= \frac{1}{\sqrt{2}} \left(\widetilde{N} + \widetilde{N}^* \right), \\ \widetilde{N}_- &= -\frac{i}{\sqrt{2}} \left(\widetilde{N} - \widetilde{N}^* \right), \end{aligned} \tag{5}$$

with the corresponding masses given by

$$M_\pm^2 = M^2 + \widetilde{M}^2 \pm BM. \tag{6}$$

In order to avoid a tachyonic mass which implies an instability of the vacuum such that the sneutrino will develop a vacuum expectation value, we always assume $B < M + \widetilde{M}^2/M$.

Rewriting the Lagrangian in terms of mass eigenstates $\widetilde{N}_\pm$ we have

$$
\begin{aligned}
-\mathcal{L}_{\widetilde{N}} - \mathcal{L}_{\rm soft} = & M_+^2 \widetilde{N}_+^* \widetilde{N}_+ + M_-^2 \widetilde{N}_-^* \widetilde{N}_- \\
& + \frac{1}{\sqrt{2}} \left\{ \widetilde{N}_+ \left[Y_\alpha \overline{\widetilde{H}_u^c} P_L \ell_\alpha + (A_\alpha + M Y_\alpha) \widetilde{\ell}_\alpha H_u \right] \right. \\
& \left. + i \widetilde{N}_- \left[Y_\alpha \overline{\widetilde{H}_u^c} P_L \ell_\alpha + (A_\alpha - M Y_\alpha) \widetilde{\ell}_\alpha H_u \right] + \text{H.c.} \right\}. \qquad (7)
\end{aligned}
$$

A. General constraints

The total decay width for $\widetilde{N}_\pm$ is given by

$$
\Gamma_\pm \simeq \frac{M}{4\pi} \sum_\alpha \left[Y_\alpha^2 + \frac{|A_\alpha|^2}{2M^2} \pm \frac{Y_\alpha \text{Re}(A_\alpha)}{M} \right], \qquad (8)
$$

where we have expanded up to $\mathcal{O}(Y_\alpha^2, m_{\rm SUSY}^2/M^2, Y_\alpha m_{\rm SUSY}/M)$ and ignored the final state phase space factors. We will impose the restriction that $|A_\alpha|, B < M$ and $Y_\alpha < 1$ to ensure that we are always in the perturbative regime. In principle, $m_{\rm SUSY}/M$ and Y_α can go up to 4π before perturbative theory breaks down but with our stronger restriction, we are not anywhere near the nonperturbative regime.

The out-of-equilibrium condition for leptogenesis is

$$
\Gamma_\pm \lesssim H(T = M), \qquad (9)
$$

where the Hubble expansion rate is given by $H = 1.66\sqrt{g_\star}\, T^2/M_{\rm Pl}$ with Planck mass $M_{\rm Pl} = 1.22 \times 10^{19}$ GeV. Assuming minimal supersymmetric Standard Model relativistic degrees of freedom, we have $g_\star = 228.75$. The condition above translates to

$$
\sqrt{\sum_\alpha \left[Y_\alpha^2 + \frac{|A_\alpha|^2}{2M^2} \pm \frac{Y_\alpha \text{Re}(A_\alpha)}{M} \right]} \lesssim 1.6 \times 10^{-5} \left(\frac{M}{10^7\,\text{GeV}} \right)^{1/2}. \qquad (10)
$$

From the condition above, we see that $|A_\alpha|$ is bounded from above depending on M. For example if $M \sim$ TeV, we require $|A_\alpha| \lesssim 10^{-4}$ GeV. At this low scale, the mass splitting between $\widetilde{N}_+$ and $\widetilde{N}_-$ is required to be of the order of their decay widths such that the CP violation is resonantly enhanced to yield successful leptogenesis [25, 26]. To avoid excessive fine-tuning, if we consider $|A_\alpha| \sim$ TeV, Eq. (10) implies $M \gtrsim 4 \times 10^7$ GeV.

In type-I seesaw, barring special cancellation, we have the upper bound on the sum of light neutrino masses from cosmology [24]

$$
\begin{aligned}
& \sum_\alpha \frac{Y_\alpha^2 v_u^2}{M} \lesssim \sum_i m_{\nu_i} \simeq 0.23\,\text{eV}, \\
& \sqrt{\sum_\alpha Y_\alpha^2} \lesssim 3 \times 10^{-4} \left(\frac{M}{10^7\,\text{GeV}} \right)^{1/2} \left(1 + \frac{1}{\tan^2\beta} \right)^{1/2}, \qquad (11)
\end{aligned}
$$

where $\tan\beta \equiv v_u/v_d$ and $v_{u(d)} = \langle H_{u(d)} \rangle$ are the up(down)-type Higgs vacuum expectation values. $v_u^2 + v_d^2 = \sqrt{2}\, G_F^{-1} \simeq (174\,\text{GeV})^2$ with G_F the Fermi constant. For $\tan\beta \gtrsim 1$, the bound above is always less stringent than Eq. (10), and, hence, the out-of-equilibrium condition alone is sufficient.

III. CP VIOLATION

In this section we will study CP violation of the Lagrangian (7) from the interferences between tree-level and one-loop diagrams shown in Figs. 1 and 3. We will take into account thermal corrections while approximating sneutrinos $\widetilde{N}_\pm$ to always be at rest with respect to the thermal bath. Since we are in the regime where all three lepton flavors can be distinguished ($T \lesssim 10^9$ GeV), we will not sum over the lepton flavor in the final states [23].

To quantify the CP violation, we define the CP asymmetry for the decays $\widetilde{N}_\pm \to a_\alpha$ with $a_\alpha = \{\widetilde{\ell}_\alpha H_u, \ell_\alpha \widetilde{H}_u\}$ as

$$
\epsilon_{\pm\alpha}^{S,V} \equiv \frac{\gamma(\widetilde{N}_\pm \to a_\alpha) - \gamma(\widetilde{N}_\pm \to \overline{a_\alpha})}{\sum_{a_\beta;\beta} \left[\gamma(\widetilde{N}_\pm \to a_\beta) + \gamma(\widetilde{N}_\pm \to \overline{a_\beta}) \right]}, \qquad (12)
$$

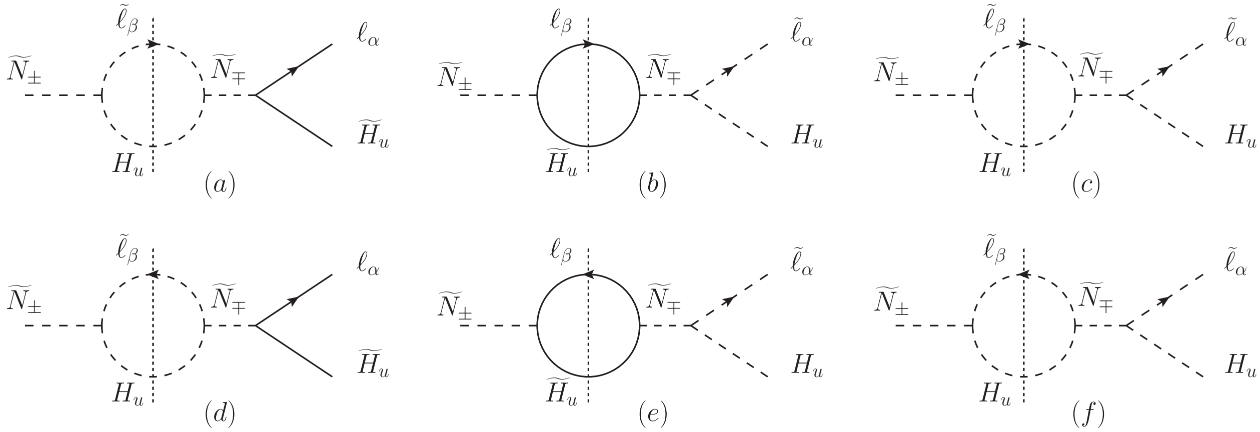

Figure 1. One-loop self-energy diagrams for the decays $\widetilde{N}_\pm \to \ell_\alpha \widetilde{H}_u$ [(a),(d)] and $\widetilde{N}_\pm \to \widetilde{\ell}_\alpha H_u$ [(b),(c),(e),(f)]. The arrow indicates the flow of lepton number. The dotted vertical lines indicate the corresponding intermediate states go on mass shell. The diagram with fermionic loop and fermionic final states does not contribute to the CP violation since it does not involve the soft couplings A_α.

where the superscripts S and V indicate the CP violation coming from self-energy and vertex corrections, respectively, $\overline{a_\alpha}$ indicates the CP conjugate of a_α, and $\gamma(i \to j)$ is the thermal averaged reaction density for the process $i \to j$ defined in Eq. (A1). In the following, we will include the thermal effects associated with intermediate on-shell states which, as shown in Ref. [20], will result in the cancellation of vertex CP asymmetries from gaugino contributions [17, 18]. We will always approximate $\widetilde{N}_\pm$ to be at rest with respect to the thermal bath so that we can obtain analytical expressions for the CP asymmetries (see Appendix A). Furthermore, we focus on the superequilibration regime which falls in the temperature range $T \lesssim 10^8$ GeV for $m_{\rm SUSY} \sim 1$ TeV [5]. The advantage is that in this regime, lepton and sleptons are not distinguished (they have the same chemical potentials) and so the two Boltzmann equations for the lepton asymmetry in particles and sparticles can be reduced to one equation for the net lepton asymmetry.[3] Hence we are allowed to sum over CP asymmetries of lepton and slepton final states as below.

A. CP violation from mixing

In this subsection, we discuss the mixing CP violation from self-energy corrections. There are two kinds of self-energy diagrams as shown in Fig. 1: the diagrams with continuous flow of lepton number [Figs. 1(a)–1(c)] and the diagrams with flow of lepton number inverted in the loop [Figs. 1(d)–1(f)]. Notice that diagrams with fermionic loop and fermionic final states do not contribute to the CP violation since they do not involve the soft couplings A_α. From Figs. 1(a)–1(c), we obtain the respective contributions to the CP asymmetries defined in Eq. (12) as follows[4]:

$$
\begin{aligned}
\epsilon^{S,(a)}_{\pm\alpha} &= \frac{1}{4\pi G_\pm(T)} Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M}\left(1+\frac{\widetilde{M}^2}{M^2} \pm \frac{B}{M}\right)\frac{2BM}{4B^2+\Gamma_\mp^2} r_B(T) c_F(T),\\
\epsilon^{S,(b)}_{\pm\alpha} &= -\frac{1}{4\pi G_\pm(T)} Y^2 Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M}\left(1+\frac{\widetilde{M}^2}{M^2} \pm \frac{B}{M}\right)\frac{2BM}{4B^2+\Gamma_\mp^2} r_F(T) c_B(T),\\
\epsilon^{S,(c)}_{\pm\alpha} &= \frac{1}{4\pi G_\pm(T)} \left[\left(Y^2 - \sum_\beta \frac{|A_\beta|^2}{M^2}\right) Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} - \left(Y_\alpha^2 - \frac{|A_\alpha|^2}{M^2}\right)\sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M}\right]\\
&\quad \times \frac{2BM}{4B^2+\Gamma_\mp^2} r_B(T) c_B(T),
\end{aligned}
\tag{13}
$$

[3] We make the assumption of superequilibration also to highlight the positive effects of nonthermal CP violation in soft leptogenesis. Including nonsuperequilibration effects, the efficiency of soft leptogensis is expected to be further enhanced, and this effect was studied in detail in Ref. [21]. The validity window of superequilibration can be enlarged by increasing the gaugino masses and μ parameter [5] and/or decreasing $|A_\alpha|$.

[4] The absorptive parts which regularize the singularity in the $\widetilde{N}_\pm$ propagators as $M_+ \to M_-$ are obtained by resumming self-energy diagrams following Refs. [25, 26].

where we define $Y^2 \equiv \sum_\alpha Y_\alpha^2$ and

$$G_\pm(T) \equiv \left[Y^2 + \sum_\alpha \left(\frac{|A_\alpha|^2}{M^2} \pm \frac{2Y_\alpha \mathrm{Re}(A_\alpha)}{M}\right)\right] c_B(T) + Y^2 \left(1 + \frac{\widetilde{M}^2}{M^2} \pm \frac{B}{M}\right) c_F(T). \tag{14}$$

In the above $r_{B,F}(T)$ and $c_{B,F}(T)$ are temperature-dependent terms associated with intermediate on-shell and final states respectively, as given in Appendix A. We will also make use of the following identity

$$r_F(T)c_B(T) = r_B(T)c_F(T), \tag{15}$$

proven in Appendix A. Note that if we sum over the lepton flavor α and use Eq. (15), we obtain $\sum_\alpha \left(\epsilon_{\pm\alpha}^{S,(a)} + \epsilon_{\pm\alpha}^{S,(b)}\right) = \sum_\alpha \epsilon_{\pm\alpha}^{S,(c)} = 0$, in agreement with the $T = 0$ result of Ref. [19] that if there is no L violation to the right of the cut in the one-loop diagrams, the net CP violation on summing over all final states is zero.

From Figs. 1(d)–1(f), we have

$$\begin{aligned}
\epsilon_{\pm\alpha}^{S,(d)} &= \frac{1}{4\pi G_\pm(T)} Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \left(1 + \frac{\widetilde{M}^2}{M^2} \pm \frac{B}{M}\right) \frac{2BM}{4B^2 + \Gamma_\mp^2} r_B(T) c_F(T), \\
\epsilon_{\pm\alpha}^{S,(e)} &= \frac{1}{4\pi G_\pm(T)} Y^2 Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} \left(1 + \frac{\widetilde{M}^2}{M^2} \pm \frac{B}{M}\right) \frac{2BM}{4B^2 + \Gamma_\mp^2} r_F(T) c_B(T), \\
\epsilon_{\pm\alpha}^{S,(f)} &= \frac{1}{4\pi G_\pm(T)} \left[-\left(Y^2 - \sum_\beta \frac{|A_\beta|^2}{M^2}\right) Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} - \left(Y_\alpha^2 - \frac{|A_\alpha|^2}{M^2}\right) \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M}\right] \\
&\quad \times \frac{2BM}{4B^2 + \Gamma_\mp^2} r_B(T) c_B(T).
\end{aligned} \tag{16}$$

Notice the leading contributions from $\widetilde{N}_+$ and $\widetilde{N}_-$ in Eqs. (13) and (16) come with the same sign and, hence, they will contribute constructively to the lepton number asymmetry.

The total CP asymmetry from mixing $\epsilon_{\pm\alpha}^S \equiv \sum_{n=\{a,b,c,d,e,f\}} \epsilon_{\pm\alpha}^{S,(n)}$ is given by

$$\begin{aligned}
\epsilon_{\pm\alpha}^S = &\frac{1}{4\pi G_\pm(T)} Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \frac{4BM}{4B^2 + \Gamma_\mp^2} \left[c_F(T) - c_B(T)\right] r_B(T) \\
&+ \frac{1}{4\pi G_\pm(T)} \frac{|A_\alpha|^2}{M^2} \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \frac{4BM}{4B^2 + \Gamma_\mp^2} r_B(T) c_B(T) \\
&+ \frac{1}{4\pi G_\pm(T)} Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \left(\frac{\widetilde{M}^2}{M^2} \pm \frac{B}{M}\right) \frac{4BM}{4B^2 + \Gamma_\mp^2} r_B(T) c_F(T).
\end{aligned} \tag{17}$$

In the above, the first term vanishes in the zero temperature limit $T \to 0$ when $c_{B,F}(T) \to 1$ and $r_{B,F}(T) \to 1$, while the terms higher order in m_{SUSY}/M survive. They remain nonzero after summing over the lepton flavor α. In the following in order to make the dependence of thermal and nonthermal CP asymmetries in Eq. (17) on the model parameters more transparent, it is instructive to look at two limiting cases (i) $Y_\alpha \gg A_\alpha/M$ and (ii) $Y_\alpha \ll A_\alpha/M$ where in case (i), the thermal CP violation dominates, while in case (ii), the nonthermal CP violation dominates.[5]

- In the limit (i) $Y_\alpha \gg A_\alpha/M$, we have

$$\epsilon_{\pm\alpha}^S \simeq \frac{1}{4\pi} P_\alpha \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \frac{4BM}{4B^2 + \Gamma_Y^2} \frac{c_F(T) - c_B(T)}{c_F(T) + c_B(T)} r_B(T), \tag{18}$$

  where we define the flavor projector $P_\alpha \equiv Y_\alpha^2/Y^2$ with $\sum_\alpha P_\alpha = 1$ and $\Gamma_Y \equiv \frac{Y^2 M}{4\pi}$, and we have dropped the terms higher order in m_{SUSY}/M. In this case, the CP asymmetry in Eq. (18) is proportional to $c_F(T) - c_B(T)$ which goes to zero as $T \to 0$, and, hence, the contribution to the CP violation is the thermal one.

[5] By thermal (nonthermal) CP violation, we refer to the case where CP violation does (not) vanish as $T \to 0$.

In the resonant regime where $B \sim \Gamma_\pm$, we have $\epsilon_\pm^S \sim (|A|/M)/Y$ where we have suppressed the lepton flavor index for an order of magnitude estimation. In this case, a large $\epsilon_\pm^S$ can be obtained which allows TeV-scale leptogenesis but at the cost of having unnaturally small $|A|, B \ll$ TeV.

Away from the resonant regime when $B \gg \Gamma_\pm$, the CP asymmetries go as $\epsilon_\pm^S \sim 10^{-1} Y|A|/B$ assuming $\mathcal{O}(1)$ contribution from the CP phases of Eq. (4). Taking $|A| \sim$ TeV $\gtrsim B$ together with the out-of-equilibrium decay condition (10) gives us sufficient CP asymmetries $\epsilon_\pm^S \gtrsim 10^{-6}$ for $M \gtrsim 10^7$ GeV.

- In the other limit (ii) $Y_\alpha \ll A_\alpha/M$, we have

$$\epsilon_{\pm\alpha}^S \simeq \frac{1}{4\pi} \frac{|A_\alpha|^2}{\sum_\delta |A_\delta|^2} \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \frac{4BM}{4B^2 + \Gamma_A^2} r_B(T), \tag{19}$$

where $\Gamma_A \equiv \sum_\alpha \frac{|A_\alpha|^2}{8\pi M}$. The CP asymmetries Eq. (19) clearly do not vanish at $T = 0$, and this represents a *nonthermal* CP violation. Of course thermal effects are always there but the fact that the CP violation is nonvanishing at $T = 0$ implies that it is less suppressed compared to case (i).

In the resonant regime $B \sim \Gamma_\pm$, we have $\epsilon_\pm^S \sim Y/(|A|/M)$. In this case too a large $\epsilon_\pm^S$ can be obtained which allows TeV-scale leptogenesis but at the cost of having unnaturally small $|A|, B \ll$ TeV.

Away from the resonant regime with $B \gg \Gamma_\pm$, the CP asymmetries, like in the limit (i), go as $\epsilon_\pm^S \sim 10^{-1} Y|A|/B$ assuming $\mathcal{O}(1)$ contribution from the CP phases of Eq. (4). Hence taking $|A| \sim$ TeV $\gtrsim B$ together with the out-of-equilibrium decay condition (10) gives us sufficient CP asymmetries $\epsilon_\pm^S \gtrsim 10^{-6}$ for $M \gtrsim 10^7$ GeV.

To confirm our estimation of successful leptogenesis and also to illustrate the enhancing effects of nonthermal CP violation, we numerically solve the Boltzmann equations using the expression for the asymmetry parameter in Eq. (17). For simplicity, we consider only decays and inverse decays of N and $\widetilde{N}_\pm$. We will also define the washout parameter as $K \equiv \Gamma_\pm/H(T = M)$ with $\Gamma_\pm$ given by Eq. (8). In Fig. 2, we plot the absolute value of the final baryon asymmetry $|Y_{\Delta B}(\infty)|$ as a function of K for the following three scenarios:

- Non Thermal dominated (NTD): In this scenario, we choose $\boldsymbol{A}/M = (10^{-4}, 10^{-2}, 1)w$ and $\boldsymbol{Y} = (10^{-5}, 10^{-3}, 10^{-1})w$.
- Thermal Dominated (TD): In this scenario, we choose $\boldsymbol{A}/M = (10^{-5}, 10^{-3}, 10^{-1})w$ and $\boldsymbol{Y} = (10^{-4}, 10^{-2}, 1)w$.
- Mixed (MIX): In this scenario, we choose $\boldsymbol{A}/M = (10^{-4}, 10^{-2}, 1)w$ and $\boldsymbol{Y} = (10^{-4}, 10^{-2}, 1)w$.

In the above $\boldsymbol{A}$ and $\boldsymbol{Y}$ represent the couplings written as 3-vectors. In all the scenarios above, we vary w between 10^{-6} and 10^{-4} such that we scan through the parameter space from the weak washout ($K = 0.1$) to the strong washout ($K = 15$) regime while still respecting the cosmological bound on the sum of light neutrino masses in Eq. (11). For definiteness, we also fix $M = 5 \times 10^7$ GeV, $\tan\beta = 10$, $\arg(A_\alpha) = -\pi/2$, and $B = 1$ TeV.

In Fig. 2, for the left plot, we solve from an initial time with $T \gg M$ assuming zero initial number densities for N and $\widetilde{N}_\pm$ while for the right plot, we assume thermal initial number densities for N and $\widetilde{N}_\pm$. For the observed baryon asymmetry, we use the recent combined Planck and WMAP CMB measurements of cosmic baryon asymmetry [24, 27] at 2σ,

$$Y_{\Delta B}^{\mathrm{CMB}} = (8.58 \pm 0.22) \times 10^{-11}\,, \tag{20}$$

which is plotted as the gray band in Fig. 2. From the plots, we see that in the NTD (blue dashed line) and MIX (purple solid line) scenarios, $|Y_{\Delta B}(\infty)|$ falls off very slowly in the strong washout regime $K > 1$. The reason is that the falloff in their efficiencies is almost completely compensated by the increases in their respective CP asymmetries as one increases A_α/M. On the other hand, in the TD (red dotted line) scenario, in the $K > 1$ regime, $|Y_{\Delta B}(\infty)|$ falls off much faster due to the additional suppression from the partial cancellation between the CP asymmetries from the decays of $\widetilde{N}_\pm$ to scalars and fermions.

In our study we are also interested in the situation when A_α and B are restricted to be around the TeV scale. In Fig. 2, the regions to the left of the thick blue dashed and purple solid vertical lines correspond to K when $A_\alpha < 5$ TeV for the NTD and MIX scenarios, respectively, while $A_\alpha < 5$ TeV for the TD scenario in the entire range of K considered in the plot. We find, for example, for the case of zero initial number densities of N and $\widetilde{N}_\pm$, the correct amount of baryon asymmetry can be obtained for NTD at $K \sim 0.8$, MIX at $K \sim 0.6$, and TD at $K \sim 4$. Notice that the appropriate sign baryon asymmetry can always be obtained by choosing the appropriate phases of the complex couplings A_α. From this numerical exercise we conclude that generation of sufficient baryon asymmetry is possible for TeV-scale A_α and $B \gg \Gamma_\pm$, i.e., far away from the resonant regime. Besides, we also see that nonthermal CP violation can significantly enhance the efficiency of soft leptogenesis.

Finally, in Appendix B we will discuss two special cases, namely, (a) $A_\alpha = AY_\alpha$ and (b) $A_\alpha = AY^2/(3Y_\alpha)$ considered in previous work.

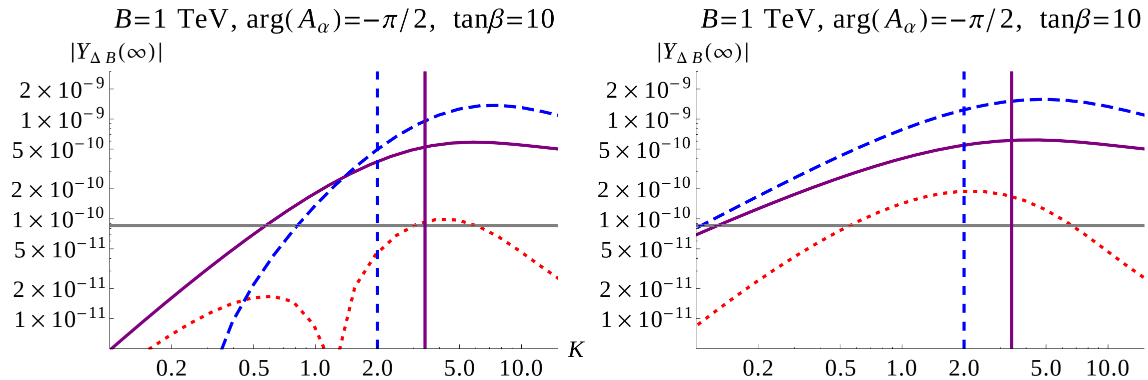

Figure 2. The absolute value of the final baryon asymmetry $|Y_{\Delta B}(\infty)|$ as a function of the washout parameter $K \equiv \Gamma_\pm/H(T = M)$ for $M = 5\times10^7$ GeV for the three scenarios described in the text: NTD (blue dashed), TD (red dotted) and MIX (purple solid). The left plot corresponds to the case of zero initial number densities of N and $\widetilde{N}_\pm$, while the right plot corresponds to the case of thermal initial number densities of N and $\widetilde{N}_\pm$. The regions to the left of the blue dashed and purple solid vertical lines correspond to K values when $A_\alpha < 5$ TeV for the NTD and MIX scenarios, respectively, while for the TD scenario we always have $A_\alpha < 5$ TeV in the range of the plot. The gray band represents the recent combined Planck and WMAP CMB measurements of cosmic baryon asymmetry [24, 27] at 2σ. The dip in the TD scenario in the left plot refers to a change in the sign of the baryon asymmetry.

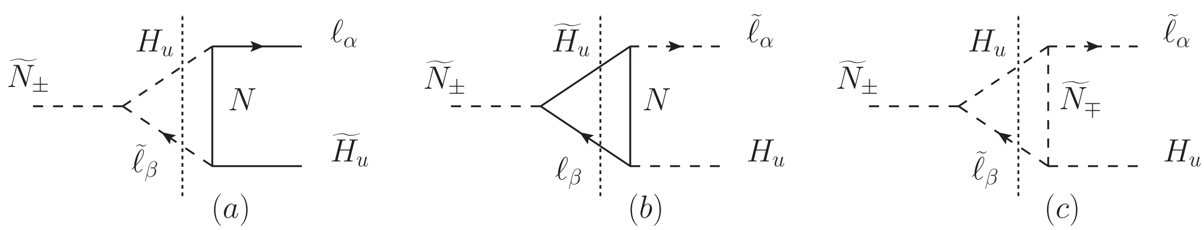

Figure 3. One-loop vertex diagrams for the decays $\widetilde{N}_\pm \to \ell_\alpha \widetilde{H}_u$ [(a)] and $\widetilde{N}_\pm \to \widetilde{\ell}_\alpha H_u$ [(b),(c)] with the conventions of Fig. 1. The diagrams with fermionic loop and fermionic final states do not contribute to the CP violation since they do not involve the soft couplings A_α.

B. CP violation from vertex corrections

In this subsection, we discuss the CP violation from vertex corrections. From Figs. 3(a)–3(c), we obtain

$$
\begin{aligned}
\epsilon^{V,(a)}_{\pm\alpha} &= \mp\frac{1}{8\pi G_\pm(T)} Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \ln\frac{M_\pm^2+M^2}{M^2} r_B(T) c_F(T), \\
\epsilon^{V,(b)}_{\pm\alpha} &= \mp\frac{1}{8\pi G_\pm(T)} Y^2 Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} \ln\frac{M_\pm^2+M^2}{M^2} r_F(T) c_B(T), \\
\epsilon^{V,(c)}_{\pm\alpha} &= \pm\frac{1}{8\pi G_\pm(T)} \left[\left(Y^2 - \sum_\beta \frac{|A_\beta|^2}{M^2}\right) Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} + \left(Y_\alpha^2 - \frac{|A_\alpha|^2}{M^2}\right)\sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M}\right] \\
&\quad \times \frac{M^2}{M_\pm^2} \ln\frac{M_\pm^2+M_\mp^2}{M_\mp^2} r_B(T) c_B(T).
\end{aligned}
\tag{21}
$$

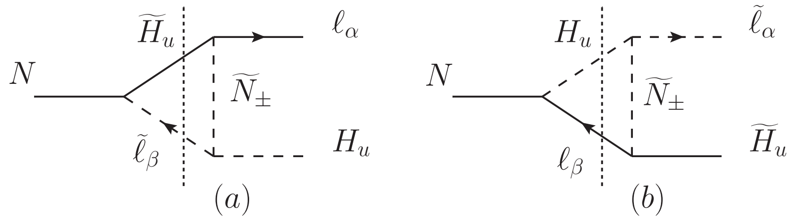

Figure 4. One-loop vertex diagrams for the decays $N \to \ell_\alpha H_u$ [(a)] and $N \to \widetilde{\ell}_\alpha \widetilde{H}_u$ [(b)] with the conventions of Fig. 1.

Summing over the contributions above and expanding in $B/M \ll 1$ in the numerators, we have

$$\begin{aligned}
\epsilon^V_{\pm\alpha} &\equiv \epsilon^{V,(a)}_{\pm\alpha} + \epsilon^{V,(b)}_{\pm\alpha} + \epsilon^{V,(c)}_{\pm\alpha} \\
&= \mp \frac{\ln 2}{8\pi G_\pm(T)} \left[Y^2 Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} + Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \right] [c_F(T) - c_B(T)]\, r_B(T) \\
&\quad - \frac{1}{8\pi G_\pm(T)} \left[Y^2 Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} + Y_\alpha^2 \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \right] \frac{B}{M} \left[\frac{c_F(T)}{2} + (\ln 2 - 1) c_B(T) \right] r_B(T) \\
&\quad \mp \frac{\ln 2}{8\pi G_\pm(T)} \left[\sum_\beta \frac{|A_\beta|^2}{M^2} Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} + \frac{|A_\alpha|^2}{M^2} \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \right] r_B(T) c_B(T) \\
&\quad + \frac{1}{8\pi G_\pm(T)} \left[\sum_\beta \frac{|A_\beta|^2}{M^2} Y_\alpha \frac{\mathrm{Im}(A_\alpha)}{M} + \frac{|A_\alpha|^2}{M^2} \sum_\beta Y_\beta \frac{\mathrm{Im}(A_\beta)}{M} \right] \frac{B}{M} (\ln 2 - 1) r_B(T) c_B(T).
\end{aligned} \tag{22}$$

The leading contributions from $\widetilde{N}_\pm$ [first and third lines of Eq. (22)] come at the order of $\epsilon^V \sim 10^{-1} Y^2$ [taking $Y_\alpha \sim \mathrm{Im}(A_\alpha)/M$], which are too small for successful leptogenesis from Eq. (10) for $M \gtrsim 10^7$ GeV. Of course the same conclusion holds also when $Y_\alpha \gg \mathrm{Im}(A_\alpha)/M$ or $Y_\alpha \ll \mathrm{Im}(A_\alpha)/M$. Besides, notice also that the leading contributions from $\widetilde{N}_\pm$ come with the opposite signs, and, hence, they will contribute destructively to the total lepton number asymmetry. Upon expanding $G_\pm(T)$ terms also in $B/M \ll 1$, we obtain an additional suppression factor B/M like the terms in the second and fourth lines in Eq. (22). Hence we conclude that the vertex CP violation is irrelevant for soft leptogenesis.

So far, we have been discussing the contributions of soft terms to CP violation in the decay of $\widetilde{N}_\pm$. In fact the soft terms also provide new sources of CP violation in the one-loop vertex diagrams for the decays of the heavy neutrino N as shown in Fig. 4. Nevertheless the CP violation from these diagrams comes at the same order as Eq. (22) and, hence, is too small for successful leptogenesis.

IV. PHENOMENOLOGICAL CONSTRAINTS

We are primarily concerned with scenarios with $M \gtrsim 10^7$ GeV for which the production of sneutrinos is beyond the energy range of current colliders. However, even if $M_\pm \sim$ TeV, the bound on the Yukawa couplings from the requirement of out-of-equilibrium decays of $\widetilde{N}_\pm$ [in Eq. (10)] makes $\widetilde{N}_\pm$ impossible to be produced at colliders [28]. On the other hand, the soft SUSY breaking parameters relevant for soft leptogenesis A_α, B, and $\widetilde{M}$ can contribute to Electric Dipole Moments (EDM) of leptons and to Charged Lepton Flavor Violating (CLFV) interactions though the analysis of Ref. [22] under the assumption of universality soft trilinear couplings $A_\alpha = AY_\alpha$ showed that the contributions to EDM and CLFV are much below the experimental bounds. Here we will repeat the analysis of Ref. [22] considering a generic A_α. In Ref. [29], the phenomenological consequences of the soft terms considering three generations of RHN chiral superfields have been discussed at length. Clearly, these soft parameters are connected with the mechanism of SUSY breaking and as such are model dependent. Here we will remain agnostic about the SUSY breaking mechanism and simply focus on the phenomenological constraints on these parameters and, in particular, we will focus only on parameters related to N_1 which are relevant for soft leptogenesis, i.e., B, $\widetilde{M}$, A_α, Y_α, and M. Without fine-tuning, we consider the soft parameters B, $\widetilde{M}$, and A_α to be similar or smaller than $m_{\mathrm{SUSY}} \sim$ TeV. On

the other hand, the parameters Y_α and M are subject only to the out-of-equilibrium $\widetilde{N}_\pm$ decay constraint in Eq. (10) and less stringently to the cosmological bound on the sum of neutrino masses in Eq. (11). The running of Y_α from the high scale down to the weak scale gives some corrections at the level of $10\% - 20\%$ (see Fig. 3 of Ref. [31]) which we will ignore in the following.

1. *Electric dipole moment of the electron*

Assuming $\mathcal{O}(1)$ contribution of the phases and mixing angles in the chargino sectors, the contributions of A_α and B to the EDM of the electron are given by [22]

$$|d_e| \approx \frac{e\, m_e \tan\beta}{16\pi m_{\tilde{\nu}}^2} \left|\frac{m_\chi Y_\alpha}{M^2}\right| \left(|A_\alpha| + BY_\alpha\right) , \tag{23}$$

where m_e is the electron mass, $m_{\tilde{\nu}}^2$ is the squared mass of the light sneutrino and m_χ the mass of chargino. For generic A_α, the first term in in Eq. (23) dominates. Taking $m_{\tilde{\nu}} = m_\chi = m_{\rm SUSY}$ and making use of Eq. (10), we have

$$|d_e| \lesssim 5 \times 10^{-38} \left(\frac{\tan\beta}{10}\right) \left(\frac{10^7\ {\rm GeV}}{M}\right)^{3/2} \left(\frac{1\ {\rm TeV}}{m_{\rm SUSY}}\right) e\,{\rm cm}, \tag{24}$$

which is much stronger than the current experimental bound $|d_e|_{\rm exp} < 8.7 \times 10^{-29} e\,{\rm cm}$ [30]. The contributions to μ and τ EDM can be estimated by replacing m_e in Eq. (23) by m_μ and m_τ, respectively, but the current experimental constraints on them are still a lot weaker: $|d_\mu|_{\rm exp} < 1.9 \times 10^{-19} e\,{\rm cm}$ [32] and $|d_\tau|_{\rm exp} < 5.1 \times 10^{-17} e\,{\rm cm}$ [33].

2. *Charged lepton flavor violating interactions*

The branching ratio for charged lepton flavor violations due to nonvanishing off-diagonal elements of the soft mass matrix of the doublet sleptons $m_{\tilde{\ell}}^2$ is given by [22, 34]

$${\rm BR}(\ell_\alpha \to \ell_\beta \gamma) \approx \frac{\alpha^3}{G_F^2} \frac{\left|(m_{\tilde{\ell}}^2)_{\alpha\beta}\right|^2}{m_{\rm SUSY}^8} \tan^2\beta, \tag{25}$$

where α is the fine structure constant. In general, the off-diagonal elements of $m_{\tilde{\ell}}^2$ will induce too-large CLFV rates. The usual solution is to assume mSUGRA boundary conditions at the grand unified theory (GUT) scale where the off-diagonal elements of $m_{\tilde{\ell}}^2$ vanish. In this case, as $m_{\tilde{\ell}}^2$ evolves from the GUT scale $M_{\rm GUT}$ to the RHN mass scale M, the off-diagonal elements will be generated due to the renormalization effects as [35]

$$(m_{\tilde{\ell}}^2)_{\alpha\beta} \approx -\frac{1}{8\pi^2} A_\alpha^* A_\beta \ln\left(\frac{M_{\rm GUT}}{M}\right) \tag{26}$$

for $\alpha \neq \beta$, and we have kept only the dominant contributions from A_α.

The most stringent constraint on the rare decay $\mu \to e\gamma$ comes from the nonobservation of the process from the MEG experiment [36, 37] which has set the new bound on the branching ratio for $\mu \to e\gamma$,

$${\rm BR}(\mu \to e\gamma)_{\rm exp} < 5.7 \times 10^{-13}. \tag{27}$$

Substituting Eq. (26) in Eq. (25) and applying the constraint Eq. (27), we obtain

$$|A_\mu^* A_e| \lesssim 5 \times 10^3\ {\rm GeV}^2 \left(\frac{m_{\rm SUSY}}{1\ {\rm TeV}}\right)^4 \left(\frac{10}{\tan\beta}\right), \tag{28}$$

where we have taken $M_{\rm GUT} = 10^{16}$ GeV and $M = 10^7$ GeV. Similarly using the experimental bounds on CLFV in τ decays, ${\rm BR}(\tau \to e\gamma)_{\rm exp} < 3.3 \times 10^{-8}$ and ${\rm BR}(\tau \to \mu\gamma)_{\rm exp} < 4.4 \times 10^{-8}$ [38], we obtain

$$|A_\tau^* A_e| \approx |A_\tau^* A_\mu| \lesssim 1 \times 10^6\ {\rm GeV}^2 \left(\frac{m_{\rm SUSY}}{1\ {\rm TeV}}\right)^4 \left(\frac{10}{\tan\beta}\right). \tag{29}$$

For $m_{\rm SUSY}$ at the TeV scale, the bound (29) can be satisfied with A_α also at the TeV scale while the stronger bound (28) requires either A_e and/or A_μ to be smaller than TeV scale. As discussed in Sec. III A, the mixing CP asymmetries away from the resonant regime go as $\epsilon_\alpha^S \sim 10^{-1} Y_\alpha |A_\alpha|/B$ and can be large enough with $M \gtrsim 10^7$ GeV and having one of the $A_\alpha \sim$ TeV $\gtrsim B$.

In addition, the off-diagonal entries of the slepton mass matrix can also give rise to other CLFV processes like $\mu \to 3e$ and $\mu - e$ conversion. If such processes are dominated by the dipole-type operator for relatively large $\tan\beta$, BR($\mu \to 3e$) and the rate of $\mu - e$ conversion rate $R_{\mu e}$ are proportional to BR($\mu \to e\gamma$) and are approximately given by [39]

$$\mathrm{BR}(\mu \to 3e) \sim 6.6 \times 10^{-3} \mathrm{BR}(\mu \to e\gamma), \tag{30}$$

and, for the ${}^{27}_{13}$Al nucleus, by [40]

$$R_{\mu e} \sim 2.5 \times 10^{-3} \mathrm{BR}(\mu \to e\gamma). \tag{31}$$

The present constraints coming from these CLFV processes are less severe than those coming from $\mu \to e\gamma$. However, in future experiments the sensitivity for such processes may improve, which could constrain the presently allowed parameter space or lead to a detection of such lepton flavor violating processes. As for example, the future Mu3e experiment [41] could reach a sensitivity of $\sim 10^{-15} - -10^{-16}$ for BR($\mu \to 3e$). For the $\mu - e$ conversion process, from the Mu2e [42] and COMET [43] experiments, the bound could reach the level of $R_{\mu e} \sim 10^{-17}$ for the ${}^{27}_{13}$Al nucleus, while the PRISM/PRIME [43] project may have 2 orders of magnitude greater sensitivity.

V. CONCLUSIONS

In the framework of local SUSY, soft leptogenesis is an attractive mechanism to explain the cosmological matter-antimatter asymmetry since it works at the lower temperature regime $T \lesssim 10^9$ GeV where the conflict with the overproduction of gravitinos can be relaxed or even evaded. We showed that by considering generic soft trilinear A_α couplings there are two interesting consequences: (1) one can realize nonthermal CP violation where the CP asymmetries in the decays of heavy sneutrinos to lepton and sleptons do not cancel at zero temperature resulting in an enhanced efficiency in generating baryon asymmetry, and (2) the dominant CP violation from self-energy corrections is sufficient even far away from the resonant regime and the relevant soft parameters can assume *natural* values at around the TeV scale. For successful soft leptogenesis, we considered two requirements: the out-of-equilibrium decays of heavy sneutrinos and a large enough CP violation. Assuming $m_{\rm SUSY} \sim$ TeV, we found the following conditions $A_\alpha \sim$ TeV $\gtrsim B$ and $M \gtrsim 10^7$ GeV as sufficient for successful leptogenesis. In addition we also found that while the contributions to the EDM of charged leptons are negligibly small, the contributions to CLFV processes are close to the sensitivities of present and future experiments.

ACKNOWLEDGMENTS

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Appendix A: Thermal corrections

The thermal averaged reaction density is defined as

$$\gamma\left(ab... \to ij...\right) \equiv \Lambda^{ij...}_{ab...} \left|\mathcal{M}\left(ab... \to ij...\right)\right|^2 f_a^{\rm eq} f_b^{\rm eq} ... \left(1 + \eta_i f_i^{\rm eq}\right) \left(1 + \eta_j f_j^{\rm eq}\right) ... , \tag{A1}$$

where $\mathcal{M}(ab... \to ij...)$ is the amplitude for the process $ab... \to ij...$ at finite temperature, $f_i^{\rm eq} = (e^{E_i/T} - \eta_i)^{-1}$ with $\eta_i = \pm$ for i representing the boson or fermion, respectively, and

$$\begin{aligned} \Lambda^{ij...}_{ab...} &\equiv \int d\Pi_a d\Pi_b ... d\Pi_i d\Pi_j ... \left(2\pi\right)^4 \delta^{(4)}\left(p_a + p_b + ... - p_i - p_j - ...\right), \\ d\Pi_i &\equiv \frac{d^3 p_i}{\left(2\pi\right)^3 2E_i}. \end{aligned} \tag{A2}$$

The CP asymmetry for the decay $a \to ij$ which arises from the interferences between tree-level and one-loop diagrams as shown in Fig. 5 is defined as

$$\begin{aligned}\epsilon &\equiv \frac{\gamma\left(a \to ij\right) - \gamma\left(a \to \overline{ij}\right)}{\sum_{k,l}\left[\gamma\left(a \to kl\right) + \gamma\left(a \to \overline{kl}\right)\right]} \\ &= \frac{\int d\Pi_a f_a^{\rm eq} \int d\Phi_{ij}\left(1+\eta_i f_i^{\rm eq}\right)\left(1+\eta_j f_j^{\rm eq}\right)\left[\left|\mathcal{M}\left(a \to ij\right)\right|^2 - \left|\mathcal{M}\left(a \to \overline{ij}\right)\right|^2\right]}{\sum_{k,l}\int d\Pi_a f_a^{\rm eq} \int d\Phi_{kl}\left(1+\eta_k f_k^{\rm eq}\right)\left(1+\eta_l f_l^{\rm eq}\right)\left[\left|\mathcal{M}\left(a \to kl\right)\right|^2 + \left|\mathcal{M}\left(a \to \overline{kl}\right)\right|^2\right]},\end{aligned} \tag{A3}$$

where the two-body phase space integral is

$$\int d\Phi_{ij} = \int d\Pi_i d\Pi_j \left(2\pi\right)^4 \delta^{(4)}\left(p_a - p_i - p_j\right). \tag{A4}$$

Ignoring the thermal motion of a with respect to the thermal bath, i.e. setting $E_a = M_a$, we can drop the integral $\int d\Pi_a$ in both the numerator and denominator. In this case the phase space integral can be carried out analytically. In order to obtain the thermal factors associated with the intermediate on-shell states, one necessarily needs to calculate the amplitudes in Eq. (A3) using thermal field theory. In the real-time formalism of thermal field theory, one needs to double the number of degrees of freedom (introducing type-1 and type-2 fields) resulting in a 2×2 matrix structure for the thermal propagator (see, e.g., [31]). However, at one loop, we can take all the vertices connected to external legs to be of type 1 and, hence, we only need to consider the (11) element of the thermal propagator. The (11) component for the boson propagator is [31]

$$D_B^{11} = \frac{i}{p^2 - m_B(T)^2 + i\epsilon} + 2\pi f_B(|p_0|)\delta(p^2 - m_B(T)^2), \tag{A5}$$

where $m_B(T)$ is the boson thermal mass, and the cut propagators are

$$D_B^{\pm} = 2\pi\left[\theta(\pm p_0) + f_B(|p_0|)\right]\delta(p^2 - m_B(T)^2). \tag{A6}$$

For fermions, the structure of the propagator is more involved [31]. For simplicity, we approximate the (11) fermion propagator by

$$D_F^{11} = \not{p}\left[\frac{i}{p^2 - m_F(T)^2 + i\epsilon} - 2\pi f_F(|p_0|)\delta(p^2 - m_F(T)^2)\right], \tag{A7}$$

where $m_F(T)$ is the fermion thermal mass and the cut propagators are

$$D_F^{\pm} = 2\pi\not{p}\left[\theta(\pm p_0) - f_F(|p_0|)\right]\delta(p^2 - m_F(T)^2). \tag{A8}$$

In Eqs. (A7) and (A8), the propagators $\sim \not{p}$ are without a mass term as the bare fermion mass is zero, and the thermal mass does not have chiral properties. Also, as implicit in the propagators, we have considered the dispersion relation as that of a free particle with a thermal mass, instead of the actual dispersion relation including thermal corrections. Although this is an underestimate of the actual dispersion relation, the error is within 10% [44]. The above also implies that in Eq. (A7) we have ignored the fact that due to the interactions with the thermal bath the two poles of the fermion propagator have different dispersion relations which can lead to an order of magnitude correction to leptogenesis in the weak washout regime and an order of 1 correction in the strong washout regime [45–47].

Keeping the above caveats in mind and applying finite temperature "cutting rules" (more discussion below), we obtain

$$\epsilon \simeq \frac{\sum_{i',j'}\left[\left|\mathcal{M}^0\left(a \to ij\right)\right|^2 - \left|\mathcal{M}^0\left(a \to \overline{ij}\right)\right|^2\right] r_{ai'j'}(T) c_{aij}(T)}{\sum_{k,l}\left[\left|\mathcal{M}^0\left(a \to kl\right)\right|^2 + \left|\mathcal{M}^0\left(a \to \overline{kl}\right)\right|^2\right] c_{akl}(T)}, \tag{A9}$$

where $\mathcal{M}^0\left(a \to ij\right)$ is the amplitude for $a \to ij$ at zero temperature and the sum over $i'j'$ in the numerator is over intermediate states in the loop which go on shell as shown as the "cuts" in Fig. 5. In Eq. (A9) the $r_{aij}(T)$'s are the thermal factors associated with the on-shell intermediate states, while the $c_{aij}(T)$'s are those associated with the final states. In the case of self-energy contributions, the factorized form as a product of thermal-dependent and zero temperature terms as in Eq. (A9) always holds (under the approximation that a is at rest with respect to the thermal

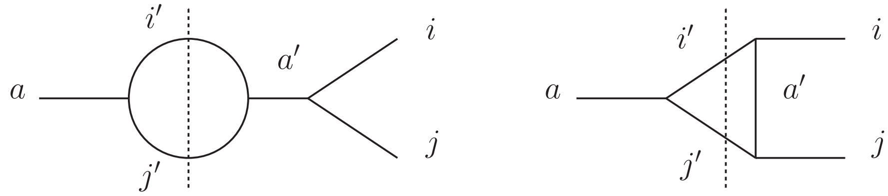

Figure 5. One-loop diagrams for decay $a \to ij$.

bath) while in the case of vertex diagrams, further approximations are required. One approximation we have made is to factorize out the temperature-dependent terms including the kinematic factors, and then to set the thermal masses in the rest of the terms to zero, which gives us expressions for these terms that coincide with the zero temperature results. In addition, we ignore the contributions from the cuts through a' and i', or a' and j' in the vertex diagrams, which as shown in Ref. [48] in non-SUSY type-I leptogenesis can give corrections depending on the $a-a'$ mass ratio, for example, at the level of 10% for $m_{a'}/m_a = 1.1$. Under these approximations, the temperature-dependent terms for both self-energy and vertex diagrams are the same and are given by

$$c_{aij}(T) = \left[1+\eta_a\left(1-\delta_{bi}\delta_{bj}\right)\left(\eta_i x_i+\eta_j x_j\right)\right]\lambda\left(1,x_i,x_j\right)\left(1+\eta_i f_i^{\rm eq}\right)\left(1+\eta_j f_j^{\rm eq}\right), \tag{A10}$$

$$r_{aij}(T) = \left[1+\eta_a\left(1-\delta_{bi}\delta_{bj}\right)\left(\eta_i x_i+\eta_j x_j\right)\right]\lambda\left(1,x_i,x_j\right)\left(1+\eta_i f_i^{\rm eq}+\eta_j f_j^{\rm eq}\right), \tag{A11}$$

with $\delta_{bi} = 1\,(0)$ if i is a boson (fermion) and

$$\lambda(1,x,y) = \sqrt{(1+x-y)^2-4x}, \qquad x_i = \frac{m_i\left(T\right)^2}{M_a^2},$$
$$E_i = \frac{M_a}{2}\left(1+x_i-x_j\right), \qquad E_j = M_a - E_i = \frac{M_a}{2}\left(1-x_i+x_j\right). \tag{A12}$$

For the statistical factors in $r_{aij}(T)$, we applied the finite temperature cutting rules by considering causal (i.e., retarded or advanced) n-point functions as pointed out by Ref. [49] which gives the dependence on the distribution functions $1+\eta_i f_i^{\rm eq}+\eta_j f_j^{\rm eq}$ in agreement with the results derived from nonequilibrium quantum field theory [48, 50–62] in contrast to the results of Refs. [31, 63] which obtained $1+\eta_i f_i^{\rm eq}+\eta_j f_j^{\rm eq}+\eta_i\eta_j f_i^{\rm eq} f_j^{\rm eq}$ when time-ordered n-point functions are considered instead. Notice that the imaginary time formalism also gives the statistical factors in agreement with the result of nonequilibrium quantum field theory [45–47].

Now we can apply the general results (A10) and (A11) to soft leptogenesis. For the decays $\widetilde{N}_\pm \to \ell_\alpha \widetilde{H}_u$ and $\widetilde{N}_\pm \to \widetilde{\ell} H_u$, the relevant thermal factors are obtained from Eqs. (A10) and (A11) to be

$$\begin{aligned}
c_F(T) &= \left(1-x_\ell-x_{\widetilde{H}_u}\right)\lambda\left(1,x_\ell,x_{\widetilde{H}_u}\right)\left(1-f_\ell^{\rm eq}\right)\left(1-f_{\widetilde{H}_u}^{\rm eq}\right),\\
c_B(T) &= \lambda\left(1,x_{\widetilde{\ell}},x_{H_u}\right)\left(1+f_{\widetilde{\ell}}^{\rm eq}\right)\left(1+f_{H_u}^{\rm eq}\right),\\
r_F(T) &= \left(1-x_\ell-x_{\widetilde{H}_u}\right)\lambda\left(1,x_\ell,x_{\widetilde{H}_u}\right)\left(1-f_\ell^{\rm eq}-f_{\widetilde{H}_u}^{\rm eq}\right),\\
r_B(T) &= \lambda\left(1,x_{\widetilde{\ell}},x_{H_u}\right)\left(1+f_{\widetilde{\ell}}^{\rm eq}+f_{H_u}^{\rm eq}\right).
\end{aligned} \tag{A13}$$

The relevant thermal masses are [31]

$$\begin{aligned}
m_{\widetilde{\ell}}(T)^2 &= 2m_\ell(T)^2 = \left(\frac{3}{8}g_2^2+\frac{1}{8}g_Y^2\right)T^2,\\
m_{H_u}(T)^2 &= 2m_{\widetilde{H}_u}(T)^2 = \left(\frac{3}{8}g_2^2+\frac{1}{8}g_Y^2+\frac{3}{4}\lambda_t^2\right)T^2.
\end{aligned} \tag{A14}$$

Next we prove a useful identity (in the context of the approximations made above) as follows

$$\begin{aligned}
r_F(T)c_B(T)-r_B(T)c_F(T) &\propto \left(1-f_\ell^{\rm eq}-f_{\widetilde{H}_u}^{\rm eq}\right)\left(1+f_{\widetilde{\ell}}^{\rm eq}\right)\left(1+f_{H_u}^{\rm eq}\right)-\left(1+f_{\widetilde{\ell}}^{\rm eq}+f_{H_u}^{\rm eq}\right)\left(1-f_\ell^{\rm eq}\right)\left(1-f_{\widetilde{H}_u}^{\rm eq}\right)\\
&= \left(e^{E_{\widetilde{N}}/T}-1\right)f_\ell^{\rm eq}f_{\widetilde{H}_u}^{\rm eq}e^{E_{\widetilde{N}}/T}f_{\widetilde{\ell}}^{\rm eq}f_{H_u}^{\rm eq}-\left(e^{E_{\widetilde{N}}/T}-1\right)f_{\widetilde{\ell}}^{\rm eq}f_{H_u}^{\rm eq}e^{E_{\widetilde{N}}/T}f_\ell^{\rm eq}f_{\widetilde{H}_u}^{\rm eq}\\
&= 0.
\end{aligned} \tag{A15}$$

In the second line above, we have made use of the following identity

$$\left(1+\eta_i f_i^{\rm eq}\right)\left(1+\eta_j f_j^{\rm eq}\right) = e^{(E_i+E_j)/T} f_i^{\rm eq} f_j^{\rm eq}, \tag{A16}$$

and the conservation of energy $E_{\widetilde{N}} = E_\ell + E_{\widetilde{H}_u} = E_{\widetilde{\ell}} + E_{H_u}$. Notice that this identity also holds if instead of using the factor $1+\eta_i f_i^{\rm eq}+\eta_j f_j^{\rm eq}$ in $r_{aij}(T)$, one uses $1+\eta_i f_i^{\rm eq}+\eta_j f_j^{\rm eq}+\eta_i\eta_j f_i^{\rm eq} f_j^{\rm eq}$ as obtained in Refs. [31, 63].

Finally, it can be shown that the CP asymmetries from gaugino contributions [17, 18] for the decays of $\widetilde{N}_\pm$ to scalars and fermions are, respectively, given by $\epsilon_g r_F(T) c_B(T)$ and $-\epsilon_g r_B(T) c_F(T)$ where ϵ_g is some temperature-independent term. Using the identity (A15), these contributions sum up to zero illustrating the cancellations pointed out by Ref. [20].

Appendix B: Special cases of mixing CP asymmetries

Here we will discuss two specific cases of mixing CP asymmetries.

(a) Universal trilinear scenario: $A_\alpha = A Y_\alpha$.

This is the scenario considered in Refs. [13, 14, 23]. In this scenario, we are always in the regime (i) $Y_\alpha \gg |A_\alpha|/M$, and from Eq. (18), we obtain [23]

$$\epsilon^S_{\pm\alpha} \simeq P_\alpha\, \bar{\epsilon}\, \frac{c_F(T)-c_B(T)}{c_F(T)+c_B(T)} r_B(T), \tag{B1}$$

with

$$\bar{\epsilon} \equiv \frac{{\rm Im}\,(A)}{M}\frac{4B\Gamma_Y}{4B^2+\Gamma_Y^2}. \tag{B2}$$

(b) Simplified misaligned scenario: $A_\alpha = A Y^2/(3Y_\alpha)$.

This is a specific scenario considered in Ref. [16]. In this scenario, we have from Eq. (14)

$$G_\pm(T) \simeq Y^2\left[c_F(T)+c_B(T)\right] + Y^2\frac{|A|^2}{M^2} d\, c_B(T), \tag{B3}$$

where $d \equiv \sum_\alpha 1/(9P_\alpha) \geq 1$, and the minimum occurs at $P_\alpha = 1/3$ for all α. In Eq. (B3), we have dropped terms of $\mathcal{O}(Y^2 m_{\rm SUSY}/M)$ except the second term which could dominate over the first when $|A|^2/M^2 \gg d^{-1}$. This condition can only be fulfilled if P_α deviates significantly from 1/3, i.e. a very hierarchical P_α.

First let us consider the case $|A|^2/M^2 \ll d^{-1}$. From Eq. (17) we obtain

$$\epsilon^S_{\pm\alpha} \simeq P_\alpha \bar{\epsilon}\, \frac{c_F\left(T\right)-c_B\left(T\right)}{c_F\left(T\right)+c_B\left(T\right)} r_B\left(T\right) + \frac{1}{9P_\alpha}\frac{\left|A\right|^2}{M^2}\bar{\epsilon}\, \frac{c_B\left(T\right)}{c_F\left(T\right)+c_B\left(T\right)} r_B\left(T\right). \tag{B4}$$

Regarding the first "thermal" term, Ref. [16] made a mistake in that the lepton flavors in the self-energy loop were not summed over resulting in expressions which were independent of P_α. Since this term coincides with Eq. (B1) there is no enhancement for the simplified misaligned scenario case compared to the universal trilinear scenario case as claimed in Ref. [16].[6]

In Ref. [16], the second "nonthermal" term is ignored assuming it is always smaller than the first one. This is always true when the P_α's are not hierarchical or only mildly hierarchical. The "nonthermal" term can dominate only in a hierarchical scenario when some of the P_α fulfill

$$P_\alpha \ll \frac{1}{3}\frac{|A|}{M}. \tag{B5}$$

[6] Nonetheless the idea of considering a generic A term to enhance the efficiency was correct as shown in the current work.

When the condition above is fulfilled, we always have a mixed scenario where we have both thermal and nonthermal contributions – since $\sum_\alpha P_\alpha = 1$, some of the P_α's cannot fulfill Eq. (B5).

For the case $|A|^2/M^2 \gg d^{-1}$, which can only happen when P_α is very hierarchical, we have from Eq. (17)

$$\epsilon^S_{\pm\alpha} \simeq P_\alpha \frac{M^2}{|A|^2 d}\,\bar{\epsilon}_A\,\frac{c_F(T) - c_B(T)}{c_B(T)} r_B(T) + \frac{1}{9P_\alpha d}\bar{\epsilon}_A\, r_B(T)\;, \tag{B6}$$

where

$$\bar{\epsilon}_A \equiv \frac{\mathrm{Im}(A)}{M}\frac{4B\Gamma_Y}{4B^2 + \Gamma_Y^2\left(\frac{|A|^2 d}{2M^2}\right)^2}. \tag{B7}$$

The condition for the second "nonthermal" term in Eq. (B6) to dominate is again Eq. (B5).

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